



LHC limits on KK-parity non-conservation in the strong sector of universal extra-dimension models



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ABSTRACT

In five-dimensional universal extra-dimensional models compactified on an S^1/Z_2 orbifold four-dimensional kinetic terms are allowed at the two fixed points. If these terms are unequal then Kaluza–Klein (KK) parity is broken. Within such a framework we consider resonant production of the $n = 1$ KK-gluon at the LHC and its subsequent decay to $t\bar{t}$, where both production and decay are KK-parity non-conserving. We use, for the first time, the exclusion data for a $t\bar{t}$ resonance obtained by the LHC experiments to limit the mass range of the lowest gluon excitation and, in a correlated fashion, of the $n = 1$ quark excitation of the KK-parity-violating model which are both found to be in the ballpark of 600–2000 GeV.

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1. Introduction

The detection of a Higgs-like scalar particle at the Large Hadron Collider (LHC) at CERN is a landmark accomplishment. Much activity is now aimed to uncover detailed properties of this new state and compare it with the expectations from the standard model (SM). There is also a continuing interest regarding the physics which lies beyond the standard model. The evidence for such physics, albeit indirect, can be traced to the issue of naturalness of the Higgs-scalar mass, the observed masses and mixing of neutrinos, and the quest for a dark matter candidate. The energy scale for new physics remains unknown but there are several motivations which encourage us to expect that it may well be within the reach of the LHC. Here our intention is to constrain a class of non-minimal universal extra-dimensional models where the lowest ($n = 1$) Kaluza–Klein excitations are not stabilized by any symmetry. We show that the data reported by the ATLAS [1,2] and the CMS [3,4] Collaborations excluding a heavy resonance in the $t\bar{t}$ channel eliminate significant regions in the $n = 1$ KK-gluon and KK-quark mass plane.

In the simplest Universal Extra Dimension (UED) model there is one extra flat spacelike dimension and it is accessible to all particles [5]. The extra dimension y is compact (radius of compactification R) and has a Z_2 symmetry ($y \rightarrow -y$) to incorporate chiral fermions. For every SM particle one has a tower of KK exci-

tations, each member being specified by an integer $n = 0, 1, 2, \dots$, the standard model particle being the $n = 0$ mode of the tower. The SM masses of the particles are small compared to $1/R$ and it is a good approximation to take the KK states for all particles at any level n to be degenerate with a mass n/R .

The Z_2 symmetry produces a conserved KK-parity which is $(-1)^n$ for the n -th KK-level. The SM particles are of even parity while those of the first level are odd. KK-parity ensures that the lightest among the $n = 1$ particles is absolutely stable and so can be a dark matter candidate, the Lightest Kaluza–Klein Particle (LKP). This is the essence of the Universal Extra Dimension (UED) Model.

The above S^1/Z_2 orbifold compactification has two fixed points at $y = 0$ and $y = \pi R$. At these boundary points inclusion of additional four-dimensional interactions is allowed by the symmetry. In fact, these terms are useful as counterterms for compensating loop-induced contributions [6] of the five-dimensional theory. In the simplest choice, the minimal Universal Extra-Dimensional Models (mUED) [7,8], these terms are chosen so that they exactly cancel the five-dimensional loop effects at the cutoff scale of the theory Λ . Thus these boundary contributions, e.g., logarithmic corrections to masses of KK particles, are such that they vanish at the scale Λ . At lower energies these contributions remove the mass degeneracy among states at the same KK-level n .

The radius of compactification R sets the mass scale of the theory and splittings within the KK-states of the same level are controlled by the cutoff Λ . They can be constrained by using known measurement results. Thus, from the muon ($g - 2$) [9], flavor changing neutral currents [10–12], $Z \rightarrow b\bar{b}$ decay [13], the

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ρ parameter [5,14], and other electroweak precision tests [15,16], one has typically $1/R \gtrsim 300\text{--}600$ GeV. Comparing the production and leptonic decay of $n = 2$ electroweak gauge bosons with the CMS LHC data a limit of $1/R \gtrsim 715$ GeV has been placed [17].

In this work we explore the scenario where the boundary terms depart from their special choice of mUED. In addition, the boundary terms at the two fixed points are allowed to be unequal. This leads to non-conservation of KK-parity and opens the door for $n = 1$ KK-states to be produced singly in SM particle collisions and also for them to decay to $n = 0$ states¹ some of which has been emphasized in an earlier work [18]. Because of these features the picture considered here is termed KK-parity violating UED.

Here we examine the coupling of the $n = 1$ KK-gluon to a pair of SM-quarks ($n = 0$ states). Such couplings are absent in mUED and are hallmarks of the non-minimality discussed above. They provide an avenue for comparing the predictions of the theory with measurements at the LHC.

Specifically, we consider resonant production of the $n = 1$ KK-gluon in pp collisions at the LHC through the KK-parity non-conserving coupling. We take the KK-gluon to be the lightest $n = 1$ level particle. Once produced, KK-conserving decays being kinematically disallowed, the KK-gluon decays to a pair of zero-mode quarks through the same KK-parity-violating coupling. The branching ratio is equal for all types of quarks and hence it is $1/6$ for the $t\bar{t}$ mode which we examine. Both the ATLAS and CMS Collaborations have looked for the signal of a resonance produced in pp collisions which decays to a pair of top-antitop quarks. Results from ATLAS for the 7 TeV [1] and 8 TeV [2] LHC runs are now available as are the corresponding findings of CMS in [3] and [4] respectively. From the lack of observation of such a state, 95% C.L. upper limits on the cross section times branching ratio of such a signal as a function of the resonance mass have been reported in these publications. Here we use the limits from the 8 TeV runs to constrain the masses of the $n = 1$ level KK quarks and gluons of the model. We would like to emphasize that this is the first effort to restrict the parameters of KK-parity-violating UED using existing LHC data.

The two essential ingredients for calculating the signal are the mass of the $n = 1$ gluon state and the strength of its KK-parity-violating couplings. In the following section we briefly review the UED scenario with boundary-localized kinetic terms and lead up to the KK-excitation masses in such a framework. In the next section we calculate the Z_2 -parity-violating coupling involving the first excitation of the KK-gluon and a zero-mode quark-antiquark pair. With these results we then derive the expected $t\bar{t}$ signal from the production of the KK-gluon at the LHC and its subsequent decay. This is compared with the CMS [4] and ATLAS [2] 8 TeV results and the restrictions on the KK-excitation masses are exhibited. Our conclusions appear at the end.

2. KK-parity-violating UED, KK masses

In nonminimal UED one can consider kinetic and mass terms localized at the fixed points. Here we restrict ourselves to boundary-localized kinetic terms only [19–24,18]. Thus we consider a five-dimensional theory with additional four-dimensional kinetic terms at the boundaries at $y = 0$ and $y = \pi R$.

We illustrate the idea by considering free fermion fields $\Psi_{L,R}$ whose zero modes are the chiral projections of the SM fermions. The five-dimensional action with BLKT is [25]

$$S = \int d^4x dy [\bar{\Psi}_L i \Gamma^M \partial_M \Psi_L + r_f^a \delta(y) \phi_L^\dagger i \bar{\sigma}^\mu \partial_\mu \phi_L$$

$$+ r_f^b \delta(y - \pi R) \phi_L^\dagger i \bar{\sigma}^\mu \partial_\mu \phi_L + \bar{\Psi}_R i \Gamma^M \partial_M \Psi_R + r_f^a \delta(y) \chi_R^\dagger i \sigma^\mu \partial_\mu \chi_R + r_f^b \delta(y - \pi R) \chi_R^\dagger i \sigma^\mu \partial_\mu \chi_R]. \quad (1)$$

Here $\sigma^\mu \equiv (I, \vec{\sigma})$ and $\bar{\sigma}^\mu \equiv (I, -\vec{\sigma})$, $\vec{\sigma}$ being the (2×2) Pauli matrices. r_f^a, r_f^b are the strengths of the boundary terms which are chosen to be the same for Ψ_L and Ψ_R for simplicity.

It is helpful to express five-dimensional fermion fields using two component chiral spinors² [25]:

$$\Psi_L(x, y) = \begin{pmatrix} \phi_L(x, y) \\ \chi_L(x, y) \end{pmatrix} = \sum_{n=0}^{\infty} \begin{pmatrix} \phi_n(x) f_L^n(y) \\ \chi_n(x) g_L^n(y) \end{pmatrix}, \quad (2)$$

$$\Psi_R(x, y) = \begin{pmatrix} \phi_R(x, y) \\ \chi_R(x, y) \end{pmatrix} = \sum_{n=0}^{\infty} \begin{pmatrix} \phi_n(x) f_R^n(y) \\ \chi_n(x) g_R^n(y) \end{pmatrix}. \quad (3)$$

Below we examine the case of Ψ_L in detail. The results for Ψ_R will be similar and can be obtained by making appropriate changes.

Variation of the action functional equation (1) utilizing Eq. (2) results in coupled equations for the y -dependent wave-functions, f_L^n, g_L^n . Following routine steps³ one finds:

$$[1 + r_f^a \delta(y) + r_f^b \delta(y - \pi R)] m_n f_L^n - \partial_y g_L^n = 0, \\ m_n g_L^n + \partial_y f_L^n = 0 \quad (n = 0, 1, 2, \dots). \quad (4)$$

Eliminating g_L^n one obtains the equations:

$$\partial_y^2 f_L^n + [1 + r_f^a \delta(y) + r_f^b \delta(y - \pi R)] m_n^2 f_L^n = 0. \quad (5)$$

From now onwards we drop the subscript L on the wave-functions and denote them simply by f and g .

The boundary conditions are [20]

$$f^n(y)|_{0-} = f^n(y)|_{0+}, \quad f^n(y)|_{\pi R+} = f^n(y)|_{\pi R-}, \quad (6)$$

$$\left. \frac{df^n}{dy} \right|_{0+} - \left. \frac{df^n}{dy} \right|_{0-} = -r_f^a m_n^2 f^n(y)|_0, \\ \left. \frac{df^n}{dy} \right|_{\pi R+} - \left. \frac{df^n}{dy} \right|_{\pi R-} = -r_f^b m_n^2 f^n(y)|_{\pi R}. \quad (7)$$

One then obtains the solutions:

$$f^n(y) = N_n \left[\cos(m_n y) - \frac{r_f^a m_n}{2} \sin(m_n y) \right], \quad 0 \leq y < \pi R, \\ f^n(y) = N_n \left[\cos(m_n y) + \frac{r_f^b m_n}{2} \sin(m_n y) \right], \quad -\pi R \leq y < 0, \quad (8)$$

where the masses m_n for $n = 0, 1, \dots$ are obtained from the transcendental equation [20]:

$$(r_f^a r_f^b m_n^2 - 4) \tan(m_n \pi R) = 2(r_f^a + r_f^b) m_n. \quad (9)$$

The solutions are *orthonormal*, i.e.:

$$\int dy [1 + r_f^a \delta(y) + r_f^b \delta(y - \pi R)] f^n(y) f^m(y) = \delta^{nm}. \quad (10)$$

These wave-functions are combinations of a sine and a cosine function unlike in the case of mUED where they are one or the other of these two trigonometric functions. This difference and that the KK masses are solutions of Eq. (9) rather than just n/R are at the root of the distinguishing features of this model.

² The Dirac gamma matrices are in the chiral representation with $\gamma_5 = \text{diag}(-I, I)$.

³ More details in the same notations and conventions can be found in [18].

¹ This is similar to R-parity violating interactions in supersymmetry.

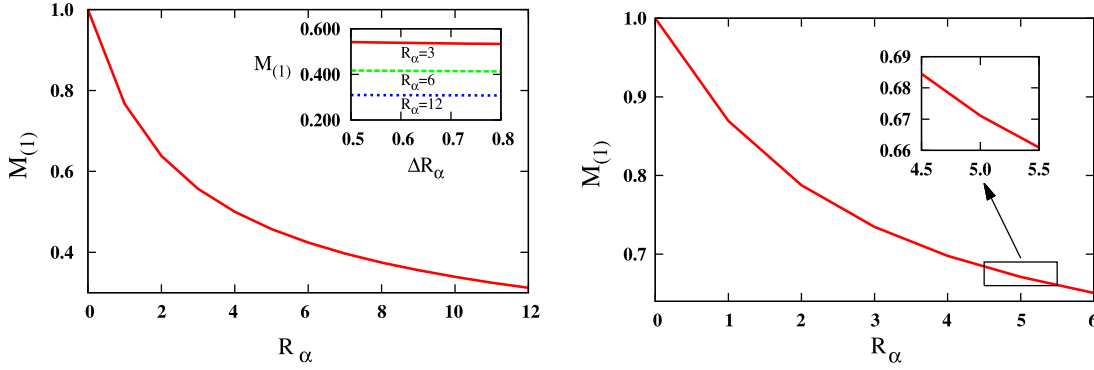


Fig. 1. Left panel: $M_{(1)} \equiv m_{\alpha(1)} R$ as a function of $R_\alpha \equiv r_\alpha^a / R$ when $r_\alpha^a = r_\alpha^b$. In the inset is shown the dependence of $M_{(1)}$ on $\Delta R_\alpha \equiv (r_\alpha^b - r_\alpha^a) / R$ for several R_α . Right panel: $M_{(1)}$ as a function of R_α when the BLKT is present only at the $y=0$ fixed point. The inset shows a blow-up of the region of R_α that is considered later. Note the very mild variation of $M_{(1)}$ in the insets of both panels. The plots apply for $\alpha = f$ (fermions) and g (gluons).

In this Letter we examine two versions of KK-parity-violating UED. In one we take symmetric boundary-localized terms for fermions, i.e., $r_f^a = r_f^b \equiv r_f$. The other case has the BLKT at one of the fixed points only: $r_f^a \neq 0, r_f^b = 0$. In this second alternative Eq. (9) becomes

$$\tan(m_n \pi R) = -\frac{r_f^a m_n}{2}. \quad (11)$$

The normalization constant N_n is determined from orthonormality. When $r_f^a = r_f^b \equiv r_f$

$$N_n = \sqrt{\frac{2}{\pi R}} \left[\frac{1}{\sqrt{1 + \frac{r_f^2 m_n^2}{4} + \frac{r_f}{\pi R}}} \right]. \quad (12)$$

For the other case when $r_f^b = 0$ we use $r_f^a \equiv r_f$ and one has

$$N_n = \sqrt{\frac{2}{\pi R}} \left[\frac{1}{\sqrt{1 + \frac{r_f^2 m_n^2}{4} + \frac{r_f}{2\pi R}}} \right]. \quad (13)$$

In this work we deal only with the zero modes and the $n=1$ KK wave-functions of the five-dimensional fermion fields.

The action for the five-dimensional gluon field, G_N ($N=0 \dots 4$) with BLKT r_g^a, r_g^b at the fixed points can be similarly written down. It is straightforward to show following similar steps⁴ that in the $G_4=0$ gauge the gluon field has the KK-expansion:

$$G_\mu(x, y) = \sum_{n=0}^{\infty} G_\mu^n(x) a^n(y), \quad (14)$$

where the functions $a^n(y)$ are of the same form as Eq. (8). In this case the five-dimensional contributions to the KK-gluon mass, m_n , satisfy

$$(r_g^a r_g^b m_n^2 - 4) \tan(m_n \pi R) = 2(r_g^a + r_g^b) m_n \quad (15)$$

which is identical to Eq. (9) for fermions.

For the other case that we also consider ($r_g^a \neq 0, r_g^b = 0$) the transcendental equation (15) reduces to

$$\tan(m_n \pi R) = -\frac{r_g^a m_n}{2}. \quad (16)$$

This equation is the same as Eq. (11) for fermions with similar BLKT.

As the KK-masses of fermions and gauge bosons are obtained from similar equations it is convenient to discuss the solutions together. Below we use r_α^a, r_α^b to stand for the BLKT strengths with $\alpha = f$ or g . Our focus will be on the $n=1$ state.

In Fig. 1 we show the variation of $M_{(1)} \equiv m_{\alpha(1)} R$, which is dimensionless, in the two alternate cases. In the left panel $M_{(1)}$ is shown as a function of $R_\alpha \equiv r_\alpha^a / R$ in the symmetric ($r_\alpha^a = r_\alpha^b$) limit. In the inset is presented the dependence of $M_{(1)}$ on the asymmetry parameter ΔR_α , for several choices of R_α . The range of variation of ΔR_α shown in the inset is what we use in our later discussion. It is seen that $M_{(1)}$ changes very slowly. Note also that the mass of the $n=1$ state for a particular R_α always remains more than that corresponding to any larger R_α for the entire variation of ΔR_α . Therefore, irrespective of the value of ΔR_α , the mass ordering within the $n=1$ level is determined on the basis of R_α . In the right panel, we present $M_{(1)}$ as a function of $R_\alpha \equiv r_\alpha^a / R$ when the BLKT is present only at $y=0$. In the inset is shown the region in which we later choose the gauge boson BLKT. Here too the variation of $M_{(1)}$ is hardly significant. The key message from both panels is that the KK-mass falls with increasing R_α ; the fermion or gauge boson with the largest R_α is the lightest $n=1$ KK state. In this work we keep $R_g > R_f$ to ensure that the $n=1$ gluon is lighter than the quarks of the same level. So, the former cannot decay via KK-number conserving modes (which would have dominated if allowed) and the branching ratio to $t\bar{t}$ is $1/6$.

3. Coupling of G^1 with zero-mode quarks

Besides the masses, the other ingredient required for the proposed calculation is the strength of the $G^1 f^0 f^0$ coupling. It is given by

$$\begin{aligned} g_{G^1 f^0 f^0} &= g_5(G) \int_0^{\pi R} (1 + r_f \{\delta(y) + \delta(y - \pi R)\}) f_L^0 f_L^0 a^1 dy \\ &= g_5(G) \int_0^{\pi R} (1 + r_f \{\delta(y) + \delta(y - \pi R)\}) g_R^0 g_R^0 a^1 dy. \end{aligned} \quad (17)$$

The five-dimensional gauge coupling g_5 which appears above is related to the usual coupling g through

$$g_5 = g \sqrt{\pi R S_G} \quad (18)$$

with

$$S_G = \left(1 + \frac{R_g^a + R_g^b}{2\pi} \right). \quad (19)$$

⁴ These steps are discussed in detail in [18].

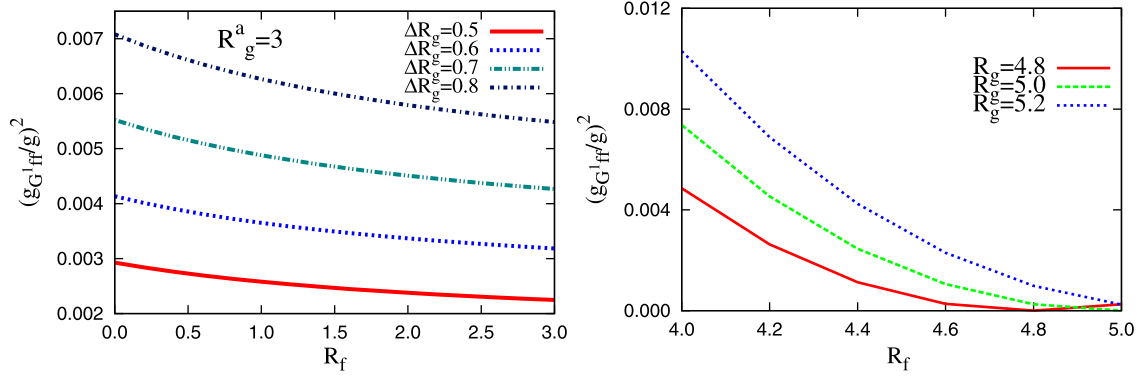


Fig. 2. Left panel: Variation of the scaled KK-parity-violating coupling squared between G^1 and a pair of zero-mode quarks with $R_f \equiv R_f^a = R_f^b$ for several ΔR_g , for $R_g^a = 3.0$. Right panel: Variation of the same coupling with R_f for different choices of R_g when the quark and boson BLKTs are present only at the $y = 0$ fixed point.

As noted in the previous section, the wave-functions f_L^0 , g_R^0 for zero-mode quarks and a^1 for the KK-gluon depend on the choices made for the boundary localized terms.

We remark in passing that, irrespective of the nature of the gluon boundary terms, the coupling $g_{G^1 G^0 G^0}$ is always zero. Thus the resonant production of G^1 is initiated only by quarks and antiquarks in the colliding particles and the gluonic content of the proton plays no role.

3.1. BLKT at both fixed points

The first option which we study has BLKTs of the same strength at the two fixed points for quarks ($r_f^a = r_f^b = r_f$) while for the gauge bosons ($r_g^a \neq r_g^b$). The y -dependent wave-functions of our interest here are found to be

$$f_L^0 = g_R^0 = \frac{1}{\sqrt{\pi R(1 + R_f/\pi)}}, \quad (20)$$

and

$$a^1 = N_G^1 \left[\cos\left(\frac{M_{(1)} y}{R}\right) - \frac{R_g^a M_{(1)}}{2} \sin\left(\frac{M_{(1)} y}{R}\right) \right], \quad (21)$$

with

$$N_G^1 = \sqrt{\frac{1}{\pi R}} \sqrt{\frac{8(4 + M_{(1)}^2 R_g^b)}{2(\frac{R_g^a + R_g^b}{\pi})(4 + M_{(1)}^2 R_g^a R_g^b) + (4 + M_{(1)}^2 R_g^a)(4 + M_{(1)}^2 R_g^b)}}, \quad (22)$$

where we have used as earlier $M_{(1)} \equiv m_{g(1)} R$, and the scaled dimensionless variables

$$R_f \equiv r_f/R, \quad R_g^a \equiv r_g^a/R, \quad \text{and} \quad R_g^b \equiv r_g^b/R. \quad (23)$$

Using the above we get

$$g_{G^1 f^0 f^0} = \frac{g(G) \sqrt{\pi R S_G}}{(1 + \frac{R_f}{\pi})} N_G^1 \left[\frac{\sin(\pi M_{(1)})}{\pi M_{(1)}} \left\{ 1 - \frac{M_{(1)}^2 R_g^a R_f}{4} \right\} + \frac{R_g^a}{2\pi} \{ \cos(\pi M_{(1)}) - 1 \} + \frac{R_f}{2\pi} \{ \cos(\pi M_{(1)}) + 1 \} \right]. \quad (24)$$

In the left panel of Fig. 2 we plot the square of the coupling for a fixed value⁵ of $R_g^a = 3.0$ as a function of R_f for several values

of ΔR_g . It is seen that the strength of the coupling decreases as R_f increases while it increases as ΔR_g increases.

Physics consequences of these couplings are discussed in the next section. At this stage we urge the reader to note that the KK-parity-violating coupling gets smaller as R_f tends towards R_g i.e., as the splitting among the $n = 1$ KK-excitations is decreased. Also, it can be readily seen using Eq. (15) that if $R_g^a = R_g^b$, i.e., the BLKTs are symmetric at $y = 0$ and $y = \pi R$ for the gauge boson, as chosen for the quarks, the coupling in Eq. (24) vanishes. This can be traced to a $y \leftrightarrow (y - \pi R)$ Z_2 -symmetry of the theory for this choice which forbids an $n = 1$ state to couple exclusively to zero modes. In general, $g_{G^1 f^0 f^0}$ decreases as ΔR_g gets smaller.

3.2. BLKT at one fixed point

In the second case, for both the quarks and the gauge bosons we assume that the BLKT are present at only the $y = 0$ fixed point. The y -dependent wave-functions in this case are

$$f_L^0 = g_R^0 = \frac{1}{\sqrt{\pi R(1 + R_f/2\pi)}}, \quad (25)$$

and

$$a^1 = \sqrt{\frac{1}{\pi R}} \sqrt{\frac{2}{1 + (\frac{R_g M_{(1)}}{2})^2 + \frac{R_g}{2\pi}}} \times \left[\cos\left(\frac{M_{(1)} y}{R}\right) - \frac{R_g M_{(1)}}{2} \sin\left(\frac{M_{(1)} y}{R}\right) \right], \quad (26)$$

where $R_g \equiv r_g/R$. With these wave-functions we get for this case

$$g_{G^1 f^0 f^0} = \frac{\sqrt{2} g(G) \sqrt{S_G}}{(1 + \frac{R_f}{2\pi}) \sqrt{1 + (\frac{R_g M_{(1)}}{2})^2 + \frac{R_g}{2\pi}}} \left(\frac{R_f - R_g}{2\pi} \right). \quad (27)$$

In the right panel of Fig. 2 we plot the square of the coupling strength as a function of R_f for several choices of R_g for this alternative. In order to keep the $n = 1$ KK-gluon lighter than the quarks of the same level we have kept $R_g > R_f$. It is to be noted that the strength of the coupling decreases as R_f increases but it increases as R_g increases. The coupling vanishes if $R_g = R_f$, as can be seen from Eq. (27).

4. Single G^1 production and decay to $t\bar{t}$

From here onwards we will not explicitly write the KK-number ($n = 0$) as a superscript for the SM particles. Wave-functions with no superscripts are for the SM particles.

⁵ We have checked that the results are not dramatically different for the other values of R_g^a that we consider later.

We now investigate the single production of G^1 at the LHC. We consider the process $pp(q\bar{q}) \rightarrow G^1$ followed by $G^1 \rightarrow t\bar{t}$. Note that both the production and decay of G^1 involves the KK-parity-violating coupling. The signature of this mode would be a resonance in the $t\bar{t}$ channel at the G^1 mass. The ATLAS and CMS Collaboration have both searched for such a resonance in pp collisions at 7 and 8 TeV [1–4]. From non-observation of such a signal 95% C.L. upper bounds have been placed on the cross section times branching ratio as a function of the mass of a $t\bar{t}$ resonance. Comparing these bounds with the calculated values in the KK-parity-violating framework enables the restriction of the parameter space of the model. To get the most up-to-date bounds we use the latest 8 TeV results. At this energy CMS has published [4] the analysis of 19.7 fb^{-1} of data while ATLAS has presented [2] bounds from 14.3 fb^{-1} of data. We use the former in our considerations below but also remark on the constraint following from the latter.

The key quantities here are the KK-gluon mass and its coupling to zero-mode quarks. In non-minimal UED, the mass of G^1 is determined from Eqs. (15) and (16) by $1/R$ and the gluon BLKT $R_g^{a,b}$. The resonance masses excluded by the ATLAS and CMS results are bounds on the $n=1$ gluon mass in this model. This restricts $1/R$ and $R_g^{a,b}$. Further, the single production and the decay of $n=1$ KK-excitations of gluons to SM quarks are driven by KK-parity-violating couplings which depend on the quark BLKT R_f and also vanish unless the strengths of the gauge BLKT parameters localized at the two fixed points are different, i.e., $\Delta R_g \neq 0$. A specific upper bound on the event rate as quoted by CMS [4] therefore translates to constraints on the above parameters and thence to the mass of the KK-excitations of quarks.

As noted earlier, in spite of the onset of KK-parity violation the coupling $g_{G^1 GG}$ vanishes identically. Consequently, the production of G^1 in pp collisions is driven solely by $q\bar{q}$ fusion. An analytic expression for the resonant production cross section of the $n=1$ KK-gluon from $q\bar{q}$ fusion in the collision of two protons can be expressed in a compact form:

$$\begin{aligned} \sigma(pp \rightarrow G^1 + X) &= \frac{4\pi^2}{3m_{g^{(1)}}^3} \sum_i \Gamma(G^1 \rightarrow q_i \bar{q}_i) \\ &\times \int_{\tau}^1 \frac{dx}{x} [f_{q_i}(\frac{x}{\tau}, m_{g^{(1)}}^2) f_{\bar{q}_i}(\frac{\tau}{x}, m_{g^{(1)}}^2) + q_i \leftrightarrow \bar{q}_i]. \end{aligned} \quad (28)$$

Here, q_i and \bar{q}_i stand for a generic quark and the corresponding antiquark of the i -th flavor respectively. $f_{q_i}(\frac{\tau}{x}, m_{g^{(1)}}^2)$ is the quark (antiquark) distribution function within a proton. $\tau \equiv m_{g^{(1)}}^2/S_{pp}$, where $\sqrt{S_{pp}}$ is the proton–proton centre of momentum energy. $\Gamma(G^1 \rightarrow q_i \bar{q}_i)$ represents the decay width of G^1 into the quark–antiquark pair and is given by

$$\Gamma = \left[\frac{g_{G^1 qq}^2}{\pi} \right] m_{g^{(1)}}. \quad (29)$$

Here $g_{G^1 qq}$ is the KK-parity-violating coupling of the $n=1$ gluon with the SM quarks – see Eqs. (24) and (27).

Eq. (28) represents the lowest order result in QCD. We have not considered higher order contributions in our analysis and used it bearing in mind that QCD corrections usually enhance cross sections and so our results are probably conservative.

To obtain the numerical values of the cross sections, we use a parton-level Monte Carlo code with parton distribution functions as parametrized in CTEQ6L [26]. We take the pp centre of momentum energy to be 8 TeV. Renormalization (for α_s) and factorization scales (in the parton distributions) are both set at $m_{g^{(1)}}$.

We are now ready to present the results of our investigation. To our knowledge, this is the first of its kind where experimental data from the LHC have been used to restrict the parameter space of KK-parity-violating UED. Results for two distinct cases, either BLKTs are present at both fixed points or only at one of the two, will be presented in following two subsections.

4.1. BLKT at both fixed points

In Fig. 3 we show the region of parameter space excluded by the CMS 8 TeV data [4] for the case in which the quark BLKTs are equal at the two fixed points but KK-parity is broken by the unequal values of the gauge BLKTs. The three panels correspond to different choices of R_g^a . In each panel the region to the left of a curve in the $m_{g^{(1)}}-R_f$ plane is disfavored by the CMS data. The curves in any one panel correspond to different choices of ΔR_g as indicated.

Since the KK-mass is rather insensitive to ΔR_g , for a chosen R_g there is one-to-one correspondence of $m_{g^{(1)}}$ with $1/R$ which is shown on the upper axis of the panels. Also, R_f determines $M_{f^{(1)}} = m_{f^{(1)}} R$ and is displayed on the right-side axis from where $m_{f^{(1)}}$ corresponding to any $1/R$ can be read off.

The results depicted in Fig. 3 can be readily understood by noting that the LHC exclusion plots translate in this model to a limit on the KK-parity violating coupling for any chosen $n=1$ gluon mass. R_g^a is fixed for a panel. In any panel, the 95% C.L. CMS upper limit on the cross section times the branching ratio implies that the points of intersection of the curves with a vertical line corresponding to a fixed $m_{g^{(1)}}$ identify $(R_f, \Delta R_g)$ pairs which lead to the same magnitude of the coupling constant. From the left panel of Fig. 2 it is seen that this happens if an increase in R_f is matched by an increase in ΔR_g . This is indeed seen to be the case. One can also easily explain the nature of the plots from the standpoint of a fixed R_f . As one moves from left to right keeping R_f fixed, i.e., to increasing $m_{g^{(1)}}$, the production cross section falls. It can be checked that though the CMS data also decreases with increasing resonance mass this is slower than the kinematic reduction. Therefore in order to match the observed results a larger KK-violating coupling is needed as $m_{g^{(1)}}$ increases. It is seen from the left panel of Fig. 2 that for a fixed R_f the coupling grows as ΔR_g is increased. This feature agrees with the results in Fig. 3.

The implications of the above results on the $n=1$ level KK mass spectrum can be extracted by considering them in conjunction with Figs. 1 and 2. The limits on $m_{g^{(1)}}$ are essentially those on the $t\bar{t}$ resonance given in the data [4]. However, for any chosen R_g^a (any one panel) this entire range cannot be covered. This is because the KK-parity violating coupling, which serves to match the model prediction for the cross section for a particular $m_{g^{(1)}}$ with the data, varies only over a limited range (see left panel of Fig. 2). This therefore determines the KK-gluon mass band permissible for a certain R_g^a . At the same time this puts an upper bound on the $n=1$ quark mass, which anyway has to be heavier than the $n=1$ gluon in this model. Thus the quark KK excitation mass has to be in a limited range to agree with the LHC data. This feature can be illustrated by a few examples from Fig. 3. From the left panel ($R_g^a = 12.0$) one finds that if $m_{g^{(1)}} = 625 \text{ GeV}$ then $m_{q^{(1)}}$ is bounded from above by 1.77 TeV. If, on the other hand $m_{g^{(1)}} = 1.60 \text{ TeV}$ the $n=1$ quark is constrained (see the right panel, $R_g^a = 3.0$) to lie between 1.60 and 2.32 TeV. In Fig. 3 the three R_g^a choices cover the entire CMS exclusion range of $t\bar{t}$ resonance mass.

4.2. BLKT at one fixed point

Now let us turn to the case of quark and gauge BLKTs at only one fixed point. In the left (right) panel of Fig. 4 we show the

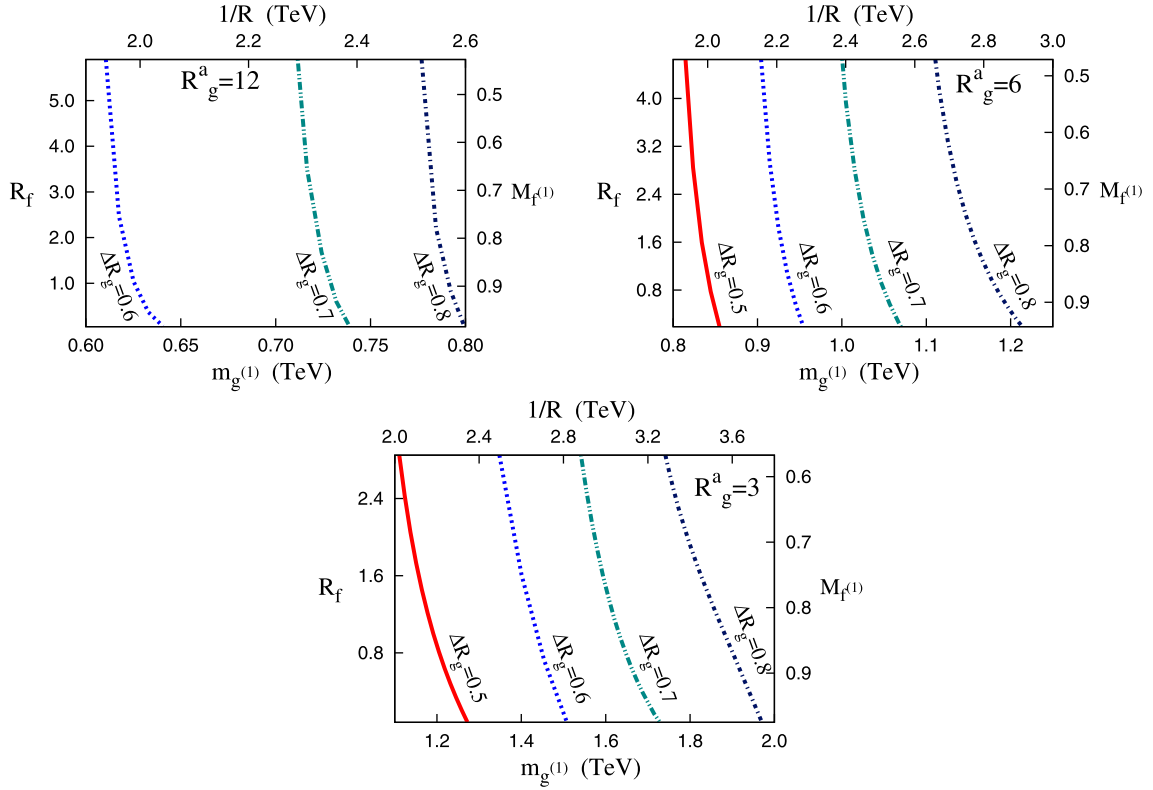


Fig. 3. 95% C.L. excluded/allowed regions in the $m_{g^{(1)}}-R_f$ parameter space for several choices of ΔR_g from non-observation of a resonant $t\bar{t}$ signal at the LHC running at 8 TeV. Each panel corresponds to a specific value of R_g^a . The region to the left of a given curve is disfavored by the CMS data [4]. $1/R$ and $M_{f^{(1)}} = m_{f^{(1)}} R$ are displayed in the upper and right-side axes respectively (see text).

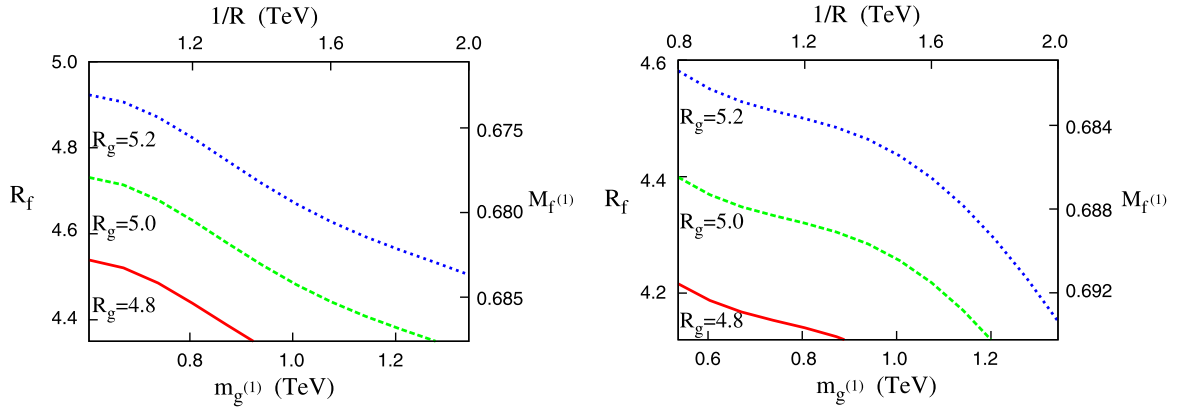


Fig. 4. 95% C.L. exclusion plots in the $m_{g^{(1)}}-R_f$ plane for several choices of R_g . The region below a specific curve is ruled out from the non-observation of a resonant $t\bar{t}$ signal in the 8 TeV run of LHC by CMS [4] (left) and ATLAS [2] (right). $1/R$ and $M_{f^{(1)}} = m_{f^{(1)}} R$ are displayed in the upper and right-side axes respectively (see text). The y-axis ranges in the two panels are different.

bounds obtained using the 8 TeV results of CMS (ATLAS). The relevant KK-parity-violating couplings vanish when⁶ $R_f = R_g$. We show in this case the exclusion curves in the $m_{g^{(1)}}-R_f$ plane for different choices of R_g . The region below a curve has been ruled out from the data. As expected the CMS bounds based on 19.7 fb^{-1} data are more restrictive than those from the 14.3 fb^{-1} ATLAS result.

Our discussion in the following is based on the left panel of Fig. 4. The nature of behavior seen in Fig. 4 tallies with our earlier considerations. As we have noted in the right panel of Fig. 1,

$M_{(1)} \equiv m_{g^{(1)}} R$ is quite insensitive to the value of R_g . It is a good approximation therefore to take the mass of G^1 to be simply proportional to $1/R$; the $1/R$ values are indicated in the upper axes of the panels in Fig. 4. For any $m_{g^{(1)}}$ the CMS data gives a bound for the corresponding cross section times branching ratio. Once the mass is fixed, the experimental limit can be achieved by a specific value for the KK-number violating coupling, i.e., the points of intersection of a vertical line with the curves give (R_g, R_f) pairs which result in this fixed coupling. This can be borne out by comparing with the right panel of Fig. 2. A second option is to consider the plots in Fig. 4 for a fixed R_f . In this case as $m_{g^{(1)}}$ increases the enhanced mass hinders the production of the $n = 1$ KK gluon. We have checked that the fall seen in the CMS data with increasing

⁶ Since for this option the BLKTs are present at only one fixed point we denote them by R_f and R_g with no superscript.

resonance mass is slower than this kinematic reduction. To compensate for this, the KK-violating coupling must increase as we move to larger $m_{g(1)}$. As seen from the right panel of Fig. 2, for a fixed R_f the coupling enhancement is accomplished by increasing R_g . This is the case in Fig. 4.

The conclusions that can be drawn from the above results are similar to the ones from Fig. 3 but much more stringent. For example, if $m_{g(1)} = 600$ GeV then depending on whether R_g is 4.8, 5.0, or 5.2 the upper bound on the $n = 1$ quark mass is 614, 610 or 605 GeV. For a heavier $n = 1$ gluon of mass 1.200 TeV one finds that the upper bound on the corresponding quark excitation is 1.236 for $R_g = 5.0$ or 1.227 TeV $R_g = 5.2$. These examples indicate that in this scenario, the $n = 1$ quarks and gluons have to be quasi-degenerate to tally with the LHC observations.

5. Conclusions

Universal extra dimension models are among the attractive options for beyond the standard model physics. Here the SM particles are complemented with heavier KK excitations which are equispaced in mass. The interaction strengths of these states are controlled entirely by the SM. Various aspects of the model ranging from constraints from precision measurements to collider searches have been looked at in the literature. Signals for UED are being actively searched for at the LHC.

One of the less attractive predictions of UED is the compressed mass spectrum of KK excitations of all SM particles at any fixed level. A remedy for this had been noted early on. It was shown [7] that five-dimensional radiative corrections split the degeneracy in a calculable way determined by the SM charges of the zero-mode states. The corrections are encoded as additional four-dimensional interactions located at the two boundary points (BLTs). In this version of UED, known as minimal UED, the practice has been to assume that the couplings of the KK excitations continue to be as for the SM particles and only the mass degeneracy is removed.

In this work we examine departures of the boundary localized kinetic terms from the above minimal choice. There are two possibilities of choosing the four-dimensional kinetic terms at the fixed points with rather distinct physics consequences. In the first, the BLKTs are of equal strength at both fixed points ($y = 0, \pi R$). Here, a Z_2 symmetry $y \longleftrightarrow (y - \pi R)$ survives. One ends up with a theory where the spectrum of KK-particles and the couplings can be drastically different from mUED. The lightest among the $n = 1$ KK particles is stable and can be a dark matter candidate [27,28]. The other alternative is to allow the BLKTs at the two fixed points to be of unequal strengths. This will lead to a breakdown of KK-parity and will allow, for example, single production of $n = 1$ KK-excitations and their decay to SM particles. Earlier we have examined, $B^1(W_3^1) \rightarrow e^+e^-, \mu^+\mu^-$, decays after the production of the $B^1(W_3^1)$ singly at the LHC [18].

In this article, we have considered the possible BLKTs for an interacting theory of quarks and gluons. In one alternative, the strengths of quark BLKTs at the two fixed points have been assumed to be equal $\equiv r_f$. For the gauge boson boundary kinetic terms we have considered the general case of unequal BLKTs ($r_g^a \neq r_g^b$). Equality of the latter strengths would restore a Z_2 -parity. As an alternate possibility we have considered the situation where the quark and gluon BLKTs are present *only* at the $y = 0$ fixed point. In both cases the boundary terms modify the field equations in the y -direction. Consistency conditions of the solutions of the above equations lead to the masses of KK-excitations of quarks and the gluons and their wave-functions in the y -direction.

We have calculated the coupling of G^1 , the $n = 1$ KK-excitation of the gluon, to a pair of zero-mode quarks (i.e., SM quarks) as

a function, in the first alternative, of r_f, r_g^a, r_g^b and $1/R$. In general, we have presented the coupling as a function of the scaled variable R_f for several choices of the other parameters. A similar KK-parity-violating coupling, which arises when the BLKTs are present only at $y = 0$, has also been evaluated. The coupling is a hallmark of KK-parity violation and vanishes in the $\Delta R_g = 0$ limit in the first case and for $R_f = R_g$ in the second.

These results are utilized to calculate the production of G^1 singly at the LHC and its subsequent decay to $t\bar{t}$, both production and decay being via the KK-parity-violating coupling. The predictions are compared with the results on $t\bar{t}$ resonance production signature at the LHC running at 8 TeV pp centre of momentum energy [2,4]. It is revealed that nonobservation of this signal with 19.7 fb^{-1} accumulated luminosity already disfavors a large part of the parameter space (spanned by r_f, r_g^a, r_g^b and $1/R$ in one case and r_f, r_g and $1/R$ in the other). In the models considered here the $n = 1$ gluon is lighter than the corresponding quark and the bounds on the mass of the former are the same as that on the $t\bar{t}$ resonance from the data. The cross section limits from LHC put tight *upper* bounds on the $n = 1$ quark excitation mass. In particular, while a range of a few hundred GeV is still permitted for this mass in the first scenario, in the second the $n = 1$ quarks and gluons have to be quasi-degenerate.

A similar new physics signal can also arise from other models, e.g., if there are extra Z -like bosons as in the left-Right symmetric models or those with an extra $U(1)$ symmetry. Here we have not attempted to compare the KK-parity-violating UED signals with those from such other scenarios.

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